

# Is SUSY accessible by direct dark matter detection?

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## Abstract

We performed a combined analysis of the parameter space of the Minimal Supersymmetric Standard Model (MSSM) taking into account cosmological and accelerator constraints including those from the radiative  $b \rightarrow s\gamma$  decay measured by the CLEO collaboration.

Special attention is paid to the event rate,  $R$ , of direct dark matter neutralino detection. We have found domains of the parameter space with  $R \simeq 5 - 10$  events/kg/day. This would be within the reach of current dark matter experiments. The  $b \rightarrow s\gamma$  data do not essentially reduce these large event rate domains of the MSSM parameter space.

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# 1 Introduction

Rare  $b \rightarrow s\gamma$  decay observed by the CLEO collaboration [1] with the branching ratio  $Br(b \rightarrow s\gamma) = (2.32 \pm 0.67) \cdot 10^{-4}$ , has been recognized as a stringent restriction for physics beyond the standard model (SM).

In the Minimal Supersymmetric Standard Model (MSSM) [2] this decay proceeds through 1-loop diagrams involving W-boson, charged Higgs boson, chargino, neutralino and gluino [3, 4]. Since the SM prediction is consistent with the measured branching ratio  $BR(b \rightarrow s\gamma)$  the MSSM contributions are stringently restricted. A dramatic reduction of the allowed MSSM parameter space due to the  $b \rightarrow s\gamma$  constraint was reported by many authors [5, 6]. The impact of this constraint on prospects for direct detection of the dark matter (DM) neutralino ( $\chi$ ) via elastic scattering off various nuclei has been also analyzed. In [5, 6] it was found that within popular supergravity models the detection rate becomes too small for observation if the CLEO constraint is incorporated in the analysis. In this case the upper bound on  $BR(b \rightarrow s\gamma)$  implies a stringent lower bound on the mass of the pseudoscalar Higgs boson ( $m_A$ ) of the MSSM if sparticles are heavy. It leads to strong suppression of the elastic neutralino-nucleus scattering cross section.

Scenarios with lighter sparticles and pseudoscalar Higgs have been studied in the literature as well [6]. Also here, first results for direct detection of neutralinos were pessimistic. Again, only small counting rates had been found in the domain of the MSSM parameter space satisfying the  $BR(b \rightarrow s\gamma)$  constraint.

A more comprehensive exploration of the constrained MSSM parameter space [7, 8, 9] discovered, however, that the  $b \rightarrow s\gamma$  constraint makes actually only a moderate effect on the expected event rate. It was realized that relatively light neutral Higgs bosons, leading to typically large event rates, are compatible with this constraint in contrast with the results of previous analyses. A key observation is that the charged Higgs boson contribution to  $BR(b \rightarrow s\gamma)$  can be efficiently compensated by the chargino contribution in a large domain of the MSSM parameter space. This essentially relaxes the lower bound both for charged and for neutral Higgses imposed by the CLEO result.

Recently an exciting result was obtained in [10]. A sophisticated scan of the MSSM parameter space constrained by the known experimental bounds including  $BR(b \rightarrow s\gamma)$  picked up points with an unexpectedly large detection rate of DM neutralinos with a mass around 1 TeV. For a germanium ( $^{76}\text{Ge}$ ) target an integrated detection rate was found at a level of 10 events/kg/day and even up to 100 events/kg/day for sodium iodide (NaI).

If these results are correct they have important consequences for direct DM detection. Such large allowed event rates would mean that the current DM experiments have already entered an unexplored part of the MSSM parameter space.

Certainly, before such a conclusion can be made one has to be sure that the above cited results, obtained in a specially arranged scan, are not an artifact having no relation to physics. One may suspect, for instance, a specific instability of the numerical code used in the analysis. Note, that the standard scan without special sampling of the model parameters did not produce even a single point with an event rate larger than 1 event/kg/day. Therefore, an independent search for large event rate points within the MSSM parameter space is apparently demanded.

In the present paper we carry out a systematic scan of the MSSM parameter space constrained by the known accelerator data and by the requirement that the DM neutralinos do not overclose the universe. We adopt the unification scenario [11] with a non-universal scalar mass when the soft Higgs mass parameters are not equal to the common sfermion soft mass parameter at the unification scale. In this case the Higgs and sfermion masses are not strongly correlated parameters. As discussed in [11], this minimal relaxation of the complete unification conditions allows one to avoid one of the most stringent theoretical limitations on the allowed values of the neutralino detection event rate. Other unification conditions do not make such an effect and tolerate large event rate values. Therefore, we keep these unification conditions to reduce a number of free parameters. The latter is crucial for a fine scanning of the MSSM parameter space which we are going to carry out.

The aim of the scan is to "detect" those domains in the parameter space where the event rate  $R$  of the direct DM detection approaches experimentally interesting values  $R > 1$  events/kg/day. Applying an extensive standard scan procedure we have found such domains with a detection rate of about 10 events/kg/day in  $^{73}\text{Ge}$ . We demonstrate that incorporation of the  $b \rightarrow s\gamma$  constraint leads to only a moderate effect on these domains.

The special sampling described in [10] has also been applied. We did not reproduce large event rate domains located around a value of 1 TeV for the neutralino mass, reported in the cited paper. No points with  $R > 1$  event/kg/day have been "detected" in this region of neutralino masses neither in the standard nor in the special scans.

The paper is organized as follows. In section 2 we specify the MSSM and give a list of the formulas relevant to our analysis. In section 3 we summarize the experimental inputs for our analysis and in section 4 we summarize the formulas for event rate calculations, then in section 5 discuss our numerical

procedure and results. Section 6 contains the conclusion.

## 2 Minimal Supersymmetric Standard Model

The MSSM is completely specified by the standard  $SU(3) \times SU(2) \times U(1)$  gauge couplings as well as by the low-energy superpotential and "soft" SUSY breaking terms [2]. The most general gauge invariant form of the R-parity conserving superpotential is

$$\mathcal{W} = h_E L^j E^c H_1^i \epsilon_{ij} + h_D Q^j D^c H_1^i \epsilon_{ij} + h_U Q^j U^c H_2^i \epsilon_{ij} + \mu H_1^i H_2^j \epsilon_{ij} \quad (1)$$

( $\epsilon_{12} = +1$ ). The following notations are used for the quark  $Q$  (3, 2, 1/6),  $D^c$  ( $\bar{3}$ , 1, 1/3),  $U^c$  ( $\bar{3}$ , 1, -2/3), lepton  $L$  (1, 2, -1/2),  $E^c$  (1, 1, 1) and Higgs  $H_1$  (1, 2, -1/2),  $H_2$  (1, 2, 1/2) chiral superfields with the  $SU(3)_c \times SU(2)_L \times U(1)_Y$  assignment given in brackets. Yukawa coupling constants  $h_{E,D,U}$  are matrices in the generation space, non-diagonal in the general case. For simplicity we suppressed generation indices.

In general, the "soft" SUSY breaking terms are given by [12]:

$$\begin{aligned} \mathcal{L}_{SB} = & -\frac{1}{2} \sum_A M_A \bar{\lambda}_A \lambda_A - m_{H_1}^2 |H_1|^2 - m_{H_2}^2 |H_2|^2 - m_{\tilde{Q}}^2 |\tilde{Q}|^2 \\ & - m_{\tilde{D}}^2 |\tilde{D}^c|^2 - m_{\tilde{U}}^2 |\tilde{U}^c|^2 - m_{\tilde{L}}^2 |\tilde{L}|^2 - m_{\tilde{E}}^2 |\tilde{E}^c|^2 \\ & - (h_E A_E \tilde{L}^j \tilde{E}^c H_1^i \epsilon_{ij} + h_D A_D \tilde{Q}^j \tilde{D}^c H_1^i \epsilon_{ij} \\ & + h_U A_U \tilde{Q}^j \tilde{U}^c H_2^i \epsilon_{ij} + \text{h.c}) - (B \mu H_1^i H_2^j \epsilon_{ij} + \text{h.c}) \end{aligned} \quad (2)$$

As usual,  $M_{3,2,1}$  are the masses of the  $SU(3) \times SU(2) \times U(1)$  gauginos  $\tilde{g}, \tilde{W}, \tilde{B}$  and  $m_i$  are the masses of scalar fields.  $A_L, A_D, A_U$  and  $B$  are trilinear and bilinear couplings.

Observable quantities can be calculated in terms of the gauge and the Yukawa coupling constants as well as the soft SUSY breaking parameters and the Higgs mass parameter  $\mu$  introduced in Eqs. (1),(2). Under the renormalization they depend on the energy scale  $Q$  according to the renormalization group equations (RGE).

It is a common practice to implement the grand unification (GUT) conditions at the GUT scale  $M_X$ . It allows one to reduce the number of free parameters of the MSSM. As explained in the introduction, we adopt a scenario with a non-universal Higgs mass with the following set of GUT conditions:

$$\begin{aligned} m_{\tilde{L}}(M_X) &= m_{\tilde{E}}(M_X) = m_{\tilde{Q}}(M_X) = m_{\tilde{U}}(M_X) = m_{\tilde{D}}(M_X) = m_0, \\ m_{H_1}(M_X) &= m_{H_2}(M_X), \end{aligned} \quad (3)$$

$$\begin{aligned} A_U(M_X) &= A_D(M_X) = A_L(M_X) = A_0, \\ M_i(M_X) &= m_{1/2}, \end{aligned} \quad (4)$$

$$\alpha_i(M_X) = \alpha_{GUT}, \text{ where } \alpha_1 = \frac{5}{3} \frac{g'^2}{4\pi}, \alpha_2 = \frac{g^2}{4\pi}, \alpha_3 = \frac{g_s^2}{4\pi}, \quad (5)$$

$g'$ ,  $g$  and  $g_s$  are the U(1), SU(2) and SU(3) gauge coupling constants. As seen from Eqs.(3) and (3) the Higgs mass parameters  $m_{H_{1,2}}$  are not equal to the common sfermion mass  $m_0$  at the GUT scale  $M_X$ . Accepting the GUT conditions above, we end up with the following free MSSM parameters: the common gauge coupling  $\alpha_{GUT}$ ; the matrices of the Yukawa couplings  $h_i^{ab}$ , where  $i = E, U, D$ ; soft supersymmetry breaking parameters  $m_0$ ,  $m_{1/2}$ ,  $A_0$ ,  $B$ , the Higgs field mixing parameter  $\mu$  and an additional parameter of the Higgs sector  $m_A$  being the mass of the CP-odd neutral Higgs boson. Since the masses of the third generation are much larger than masses of the first two ones, we consider only the Yukawa coupling of the third generation and drop the indices  $a, b$ .

Additional constraints follow from the minimization conditions of the scalar Higgs potential. Using these conditions the bilinear coupling  $B$  can be replaced in the given list of free parameters by the ratio  $\tan \beta = v_2/v_1$  of the vacuum expectation values of the two Higgs doublets.

We calculate the Fermi-scale parameters in Eqs.(1) and (2) in terms of the above listed free parameters on the basis of 2-loop RGEs following the iteration algorithm described in [13].

The Higgs potential  $V$  including the one-loop corrections  $\Delta V$  can be written as:

$$\begin{aligned} V(H_1^0, H_2^0) &= m_1^2 |H_1^0|^2 + m_2^2 |H_2^0|^2 - m_3^2 (H_1^0 H_2^0 + \text{h.c.}) \\ &+ \frac{g^2 + g'^2}{8} (|H_1^0|^2 - |H_2^0|^2)^2 + \Delta V, \end{aligned} \quad (6)$$

$$\text{with } \Delta V = \frac{1}{64\pi^2} \sum_i (-1)^{2J_i} (2J_i + 1) C_i m_i^4 \left[ \ln \frac{m_i^2}{Q^2} - \frac{3}{2} \right],$$

where the sum is taken over all possible particles with the spin  $J_i$  and with the color degrees of freedom  $C_i$ . The mass parameters of the potential are introduced in the usual way as

$$m_{1,2}^2 = m_{H_{1,2}}^2 + \mu^2, \quad m_3^2 = B\mu, \quad (7)$$

They are running parameters with the scale Q-dependence  $m_i(Q)$  determined by the RGE. The 1-loop potential (7) itself is Q-independent up to, field-independent term depending on Q, irrelevant for the symmetry breaking.

At the minimum of this potential the neutral components of the Higgs field acquire non-zero vacuum expectation values  $\langle H_{1,2}^0 \rangle = v_{1,2}$  triggering the electroweak symmetry breaking with  $g^2(v_1^2 + v_2^2) = 2M_W^2$ .

The minimization conditions read

$$2m_1^2 = 2m_3^2 \tan \beta - M_Z^2 \cos 2\beta - 2\Sigma_1 \quad (8)$$

$$2m_2^2 = 2m_3^2 \cot \beta + M_Z^2 \cos 2\beta - 2\Sigma_2, \quad (9)$$

where  $\Sigma_k \equiv \frac{\partial \Delta V}{\partial \psi_k}$ , with  $\psi_{1,2} = \mathbf{Re} H_{1,2}^0$ , are the one-loop corrections [14]:

$$\Sigma_k = -\frac{1}{32\pi^2} \sum_i (-1)^{2J_i} (2J_i + 1) \frac{1}{\psi_k} \frac{\partial m_i^2}{\partial \psi_k} m_i^2 \left( \log \frac{m_i^2}{Q^2} - 1 \right) \quad (10)$$

As remnant of two Higgs doublets  $H_{1,2}$  after the electroweak symmetry breaking there occur five physical Higgs particles: CP-odd neutral Higgs boson  $A$ , CP-even neutral Higgs bosons  $H, h$  and a pair of charged Higgses  $H^\pm$ . Their masses  $m_A, m_{h,H}, m_{H^\pm}$  can be calculated including all 1-loop corrections as second derivatives of the Higgs potential in Eq. (7) with respect to the corresponding fields evaluated at the minimum [15, 16].

The neutralino mass matrix written in the basis  $(\tilde{B}, \tilde{W}^3, \tilde{H}_1^0, \tilde{H}_2^0)$  has the form:

$$\mathcal{M}_\chi = \begin{pmatrix} M_1 & 0 & -M_Z c_\beta s_w & M_Z s_\beta s_w \\ 0 & M_2 & M_Z c_\beta c_w & -M_Z s_\beta c_w \\ -M_Z c_\beta s_w & M_Z c_\beta c_w & 0 & -\mu \\ M_Z s_\beta s_w & -M_Z s_\beta c_w & -\mu & 0 \end{pmatrix}, \quad (11)$$

where  $s_w = \sin \theta_W$ ,  $c_w = \cos \theta_W$  and  $s_\beta = \sin \beta$ ,  $c_\beta = \cos \beta$ . Diagonalizing the mass matrix above by virtue of the orthogonal matrix  $\mathcal{N}$  one can obtain the four physical neutralinos  $\chi_i$  with the field content

$$\chi_i = \mathcal{N}_{i1} \tilde{B} + \mathcal{N}_{i2} \tilde{W}^3 + \mathcal{N}_{i3} \tilde{H}_1^0 + \mathcal{N}_{i4} \tilde{H}_2^0. \quad (12)$$

and with masses  $m_{\chi_i}$  being eigenvalues of the mass matrix (11). The lightest neutralino  $\chi_1$  we denote  $\chi$ . In our analysis  $\chi$  is the lightest SUSY particle (LSP).

The chargino mass term is

$$(\tilde{W}^-, \tilde{H}_1^-) \mathcal{M}_{\tilde{\chi}^\pm} \begin{pmatrix} \tilde{W}^+ \\ \tilde{H}_2^+ \end{pmatrix} + \text{h.c.} \quad (13)$$

with the mass matrix

$$\mathcal{M}_{\tilde{\chi}^\pm} = \begin{pmatrix} M_2 & \sqrt{2} M_W \sin \beta \\ \sqrt{2} M_W \cos \beta & \mu \end{pmatrix} \quad (14)$$

which can be diagonalized by the transformation

$$\tilde{\chi}^- = U_{i1}\tilde{W}^- + U_{i2}\tilde{H}^-, \quad \tilde{\chi}^+ = V_{i1}\tilde{W}^+ + V_{i2}\tilde{H}^+ \quad (15)$$

with  $U^*\mathcal{M}_{\tilde{\chi}^\pm}V^\dagger = \text{diag}(M_{\tilde{\chi}_1^\pm}, M_{\tilde{\chi}_2^\pm})$ , where the chargino masses are

$$M_{\tilde{\chi}_{1,2}^\pm}^2 = \frac{1}{2} \left[ M_2^2 + \mu^2 + 2M_W^2 \mp \sqrt{(M_2^2 - \mu^2)^2 + 4M_W^4 \cos^2 2\beta + 4M_W^2(M_2^2 + \mu^2 + 2M_2\mu \sin 2\beta)} \right]$$

The mass matrices for the 3-generation sfermions  $\tilde{t}$ ,  $\tilde{b}$  and  $\tilde{\tau}$  in the  $\tilde{f}_L$ - $\tilde{f}_R$  basis are:

$$\begin{aligned} \mathcal{M}_{\tilde{t}}^2 &= \begin{pmatrix} m_{\tilde{Q}}^2 + m_t^2 + \frac{1}{6}(4M_W^2 - M_Z^2) \cos 2\beta & m_t(A_t - \mu \cot \beta) \\ m_t(A_t - \mu \cot \beta) & m_{\tilde{U}}^2 + m_t^2 - \frac{2}{3}(M_W^2 - M_Z^2) \cos 2\beta \end{pmatrix} \\ \mathcal{M}_{\tilde{b}}^2 &= \begin{pmatrix} m_{\tilde{Q}}^2 + m_b^2 - \frac{1}{6}(2M_W^2 + M_Z^2) \cos 2\beta & m_b(A_b - \mu \tan \beta) \\ m_b(A_b - \mu \tan \beta) & m_{\tilde{D}}^2 + m_b^2 + \frac{1}{3}(M_W^2 - M_Z^2) \cos 2\beta \end{pmatrix} \\ \mathcal{M}_{\tilde{\tau}}^2 &= \begin{pmatrix} m_{\tilde{L}}^2 + m_\tau^2 - \frac{1}{2}(2M_W^2 - M_Z^2) \cos 2\beta & m_\tau(A_\tau - \mu \tan \beta) \\ m_\tau(A_\tau - \mu \tan \beta) & m_{\tilde{E}}^2 + m_\tau^2 + (M_W^2 - M_Z^2) \cos 2\beta \end{pmatrix} \end{aligned}$$

For simplicity we ignored in the sfermion mass matrices a non-diagonality in the generation space which is important only for the  $b \rightarrow s\gamma$ -decay.

### 3 Constrained MSSM parameter space

In this section we shortly summarize the theoretical and experimental constraints used in our analysis.

Solution of the gauge coupling constants unification (see Eq.(5)) using 2-loop RGEs allows us to define the unification scale  $M_X$ . The following standard definitions are used:  $\alpha_1 = 5\alpha/(3 \cos^2 \theta_W)$ ,  $\alpha_2 = \alpha/\sin^2 \theta_W$ . The world averaged values of the gauge couplings at the  $Z^0$  energy were obtained from a fit to the LEP data [17],  $M_W$  [18] and  $m_t$  [19, 20]:  $\alpha^{-1}(M_Z) = 128.0 \pm 0.1$ ,  $\sin^2 \theta_{\overline{MS}} = 0.2319 \pm 0.0004$ ,  $\alpha_3 = 0.125 \pm 0.005$ . The value of  $\alpha^{-1}(M_Z)$  was updated from [21] by using new data on the hadronic vacuum polarization [22].

SUSY particles have not been found so far and from the searches at LEP one knows that the lower limit on the charged sleptons is half the  $Z^0$  mass (45 GeV) to be above 60 GeV. For the charginos the preliminary lower limit of 65 GeV from the LEP 140 GeV run was used. The lower limit on the lightest neutralino is 18.4 GeV while the sneutrinos have to be above 41 GeV.

Radiative corrections trigger spontaneous symmetry breaking in the electroweak sector. In this case the Higgs potential has its minimum for non-zero vacuum expectation values of the fields. Solving for  $M_Z$  from Eqs. (8) and (9) yields:

$$\frac{M_Z^2}{2} = \frac{m_1^2 + \Sigma_1 - (m_2^2 + \Sigma_2) \tan^2 \beta}{\tan^2 \beta - 1}, \quad (16)$$

where the  $\Sigma_1$  and  $\Sigma_2$  are defined in Eq. (10). This is an important constraint which relates the true vacuum to the physical Z-boson mass  $M_Z = 91.187 \pm 0.007 \text{ GeV}$ .

Another stringent constraint is imposed by the branching ratio  $BR(b \rightarrow s\gamma)$  measured by the CLEO collaboration [1] to be:  $BR(b \rightarrow s\gamma) = (2.32 \pm 0.67) \times 10^{-4}$ .

In the MSSM this flavour changing neutral current (FCNC) receives in addition to the SM  $W - t$  loop contributions from  $H^\pm - t$ ,  $\tilde{\chi}^\pm - \tilde{t}$  and  $\tilde{g} - \tilde{q}$  loops. The  $\chi - \tilde{t}$  loops, which are expected to be much smaller, have been neglected [3, 26]. The  $\tilde{g} - \tilde{q}$  loops are proportional to  $\tan \beta$ . It was found [13] that this contribution should be small, even in the case of large  $\tan \beta$  and therefore it was neglected. The chargino contribution, which becomes large for large  $\tan \beta$  and small chargino masses, depends sensitively on the splitting of the two stop masses.

Within the MSSM the following ratio has been calculated [3]:

$$\frac{BR(b \rightarrow s\gamma)}{BR(b \rightarrow ce\bar{\nu})} = \frac{|V_{ts}^* V_{tb}|^2}{|V_{cb}|^2} K_{NLO}^{QCD} \frac{6\alpha}{\pi} \frac{[\eta^{16/23} A_\gamma + \frac{8}{3}(\eta^{14/23} - \eta^{16/23}) A_g + C]^2}{I(m_c/m_b)[1 - (2/3\pi)\alpha_s(m_b)f(m_c/m_b)]}, \quad (17)$$

where

$$C \approx 0.175, \quad I = 0.4847, \quad \eta = \alpha_s(M_W)/\alpha_s(m_b), \quad f(m_c/m_b) = 2.41.$$

Here  $f(m_c/m_b)$  represents corrections from leading order QCD to the known semileptonic  $b \rightarrow ce\bar{\nu}$  decay rate, while the ratio of masses of  $c$ - and  $b$ -quarks is taken to be  $m_c/m_b = 0.316$ . The ratio of CKM matrix elements  $\frac{|V_{ts}^* V_{tb}|^2}{|V_{cb}|^2} = 0.95$  was taken from [27], the next leading order QCD-corrections from [28].  $A_{\gamma,g}$  are the coefficients of the effective operators for  $bs\text{-}\gamma$  and for  $bs\text{-}g$  interactions respectively.

Assuming that the neutralinos form a dominant part of the DM in the universe one obtains a cosmological constraint on the neutralino relic density.



The present lifetime of the universe is at least  $10^{10}$  years, which implies an upper limit on the expansion rate and correspondingly on the total relic abundance. Assuming  $h_0 > 0.4$  one finds that the contribution of each relic particle species  $\chi$  has to obey

$$\Omega_\chi h_0^2 < 1, \quad (18)$$

where the relic density parameter  $\Omega_\chi = \rho_\chi/\rho_c$  is the ratio of the relic neutralino mass density  $\rho_\chi$  to the critical one  $\rho_c = 1.88 \cdot 10^{-29} \text{h}_0^2 \text{g} \cdot \text{cm}^{-3}$ .

We calculate  $\Omega_\chi h_0^2$  following the standard procedure on the basis of the approximate formula [31, 32]:

$$\Omega_\chi h_0^2 = \frac{2.13}{10^{11}} \left( \frac{T_\chi}{T_\gamma} \right)^3 \left( \frac{T_\gamma}{2.7 K^o} \right)^3 N_F^{1/2} \left( \frac{\text{GeV}^{-2}}{ax_F + bx_F^2/2} \right). \quad (19)$$

Here  $T_\gamma$  is the present day photon temperature,  $T_\chi/T_\gamma$  is the reheating factor,  $x_F = T_F/m_\chi \approx 1/20$ ,  $T_F$  is the neutralino freeze-out temperature, and  $N_F$  is the total number of degrees of freedom at  $T_F$ . The coefficients  $a, b$  are determined from the non-relativistic expansion

$$< \sigma_{\text{ann}y} v > \approx a + bx \quad (20)$$

of the thermally averaged cross section of neutralino annihilation. We adopt an approximate treatment not taking into account complications, which occur when the expansion (20) fails [30]. We take into account all possible channels of the  $\chi - \chi$  annihilation. The most complete list of the relevant formulas for the coefficients  $a, b$  and numerical values for the other parameters in eqs. (19) and (20) can be found in [32].

Since the neutralinos are mixtures of gauginos and higgsinos, the annihilation can occur both, via s-channel exchange of the  $Z^0$  and Higgs bosons and t-channel exchange of a scalar particle, like a selectron [33]. This constrains the parameter space, as discussed by many groups [7, 32, 34, 35]. The size of the Higgsino component depends on the relative sizes of the elements in the mixing matrix (11), especially on  $\tan\beta$  and the size of the parameter  $\mu$ .

In the analysis we ignore possible rescaling of the local neutralino density  $\rho$  which may occur in the region of the MSSM parameter space where  $\Omega_\chi h^2 < 0.025$  [8, 36, 37]. This is a minimal value corresponding to DM concentrated in galactic halos averaged over the universe. If the neutralino is accepted as a dominant part of the DM its density has to exceed the quoted limiting value 0.025. Otherwise the presence of additional DM components should be taken into account, for instance, by the mentioned rescaling ansatz. However, the

halo density is known to be very uncertain. Its actual value can be one order of magnitude smaller. Therefore, one can expect that the rescaling takes place in a small domain of the MSSM parameter space. Another point is that the SUSY solution of the DM problem with such low neutralino density becomes questionable. We assume neutralinos to be a dominant component of the DM halo of our galaxy with a density  $\rho_\chi = 0.3 \text{ GeV}\cdot\text{cm}^{-3}$  in the solar vicinity and disregard in the analysis points with  $\Omega_\chi h^2 < 0.025$ .

## 4 Neutralino-Nucleus Elastic Scattering

A dark matter event is elastic scattering of a DM neutralino from a target nucleus producing a nuclear recoil which can be detected by a suitable detector. The corresponding event rate depends on the distribution of the DM neutralinos in the solar vicinity and the cross section  $\sigma_{el}(\chi A)$  of neutralino-nucleus elastic scattering. In order to calculate  $\sigma_{el}(\chi A)$  one should specify neutralino-quark interactions. The relevant low-energy effective Lagrangian can be written in a general form as

$$L_{eff} = \sum_q \left( \mathcal{A}_q \cdot \bar{\chi} \gamma_\mu \gamma_5 \chi \cdot \bar{q} \gamma^\mu \gamma_5 q + \frac{m_q}{M_W} \cdot \mathcal{C}_q \cdot \bar{\chi} \chi \cdot \bar{q} q \right) + O\left(\frac{1}{m_{\tilde{q}}^4}\right), \quad (21)$$

where terms with vector and pseudoscalar quark currents are omitted being negligible in the case of non-relativistic DM neutralinos with typical velocities  $v_\chi \approx 10^{-3}c$ .

In the Lagrangian (21) we also neglect terms which appear in supersymmetric models at the order of  $1/m_{\tilde{q}}^4$  and higher, where  $m_{\tilde{q}}$  is the mass of the scalar superpartner  $\tilde{q}$  of the quark  $q$ . These terms, as recently pointed out in [32], are potentially important in the spin-independent neutralino-nucleon scattering, especially in domains of the MSSM parameter space where  $m_{\tilde{q}}$  is close to the neutralino mass  $m_\chi$ . Below we adopt the approximate treatment of these terms proposed in [32] which allows "effectively" absorbing them into the coefficients  $\mathcal{C}_q$  in a wide region of the SUSY model parameter space.

$$\begin{aligned} \mathcal{A}_q = & - \frac{g_2^2}{4M_W^2} \left[ \frac{\mathcal{N}_{14}^2 - \mathcal{N}_{13}^2}{2} T_3 \right. \\ & - \frac{M_W^2}{m_{\tilde{q}1}^2 - (m_\chi + m_q)^2} (\cos^2 \theta_q \phi_{qL}^2 + \sin^2 \theta_q \phi_{qR}^2) \\ & \left. - \frac{M_W^2}{m_{\tilde{q}2}^2 - (m_\chi + m_q)^2} (\sin^2 \theta_q \phi_{qL}^2 + \cos^2 \theta_q \phi_{qR}^2) \right] \end{aligned}$$

$$\begin{aligned}
& - \frac{m_q^2}{4} P_q^2 \left( \frac{1}{m_{\tilde{q}1}^2 - (m_\chi + m_q)^2} + \frac{1}{m_{\tilde{q}2}^2 - (m_\chi + m_q)^2} \right) \\
& - \frac{m_q}{2} M_W P_q \sin 2\theta_q T_3 (\mathcal{N}_{12} - \tan \theta_W \mathcal{N}_{11}) \\
& \times \left( \frac{1}{m_{\tilde{q}1}^2 - (m_\chi + m_q)^2} - \frac{1}{m_{\tilde{q}2}^2 - (m_\chi + m_q)^2} \right) \Big] \quad (22)
\end{aligned}$$

$$\begin{aligned}
\mathcal{C}_q = & - \frac{g_2^2}{4} \left[ \frac{F_h}{m_h^2} h_q + \frac{F_H}{m_H^2} H_q \right. \\
& + P_q \left( \frac{\cos^2 \theta_q \phi_{qL} - \sin^2 \theta_q \phi_{qR}}{m_{\tilde{q}1}^2 - (m_\chi + m_q)^2} - \frac{\cos^2 \theta_q \phi_{qR} - \sin^2 \theta_q \phi_{qL}}{m_{\tilde{q}2}^2 - (m_\chi + m_q)^2} \right) \\
& + \sin 2\theta_q \left( \frac{m_q}{4M_W} P_q^2 - \frac{M_W}{m_q} \phi_{qL} \phi_{qR} \right) \\
& \times \left. \left( \frac{1}{m_{\tilde{q}1}^2 - (m_\chi + m_q)^2} - \frac{1}{m_{\tilde{q}2}^2 - (m_\chi + m_q)^2} \right) \right]. \quad (23)
\end{aligned}$$

Here

$$\begin{aligned}
F_h &= (\mathcal{N}_{12} - \mathcal{N}_{11} \tan \theta_W) (\mathcal{N}_{14} \cos \alpha_H + \mathcal{N}_{13} \sin \alpha_H), \\
F_H &= (\mathcal{N}_{12} - \mathcal{N}_{11} \tan \theta_W) (\mathcal{N}_{14} \sin \alpha_H - \mathcal{N}_{13} \cos \alpha_H), \\
h_q &= \left( \frac{1}{2} + T_3 \right) \frac{\cos \alpha_H}{\sin \beta} - \left( \frac{1}{2} - T_3 \right) \frac{\sin \alpha_H}{\cos \beta}, \\
H_q &= \left( \frac{1}{2} + T_3 \right) \frac{\sin \alpha_H}{\sin \beta} + \left( \frac{1}{2} - T_3 \right) \frac{\cos \alpha_H}{\cos \beta}, \\
\phi_{qL} &= \mathcal{N}_{12} T_3 + \mathcal{N}_{11} (Q - T_3) \tan \theta_W, \\
\phi_{qR} &= \tan \theta_W Q \mathcal{N}_{11}, \\
P_q &= \left( \frac{1}{2} + T_3 \right) \frac{\mathcal{N}_{14}}{\sin \beta} + \left( \frac{1}{2} - T_3 \right) \frac{\mathcal{N}_{13}}{\cos \beta}
\end{aligned}$$

Our formulas for the coefficients  $\mathcal{A}_q$  and  $\mathcal{C}_q$  of the effective Lagrangian take into account squark mixing  $\tilde{q}_L - \tilde{q}_R$  and the contribution of both CP-even Higgs bosons  $h, H$ . The formulas coincide with the relevant formulas in [32] neglecting the terms  $\sim 1/m_{\tilde{q}}^4$  and higher. These terms are taken into account "effectively" by introducing an "effective" stop quark  $\tilde{t}$  propagator.

A general representation of the differential cross section of neutralino-nucleus scattering can be given in terms of three spin-dependent  $\mathcal{F}_{ij}(q^2)$  and one spin-independent  $\mathcal{F}_S(q^2)$  form factors as follows [38]

$$\frac{d\sigma}{dq^2}(v, q^2) = \frac{8G_F}{v^2} \left( a_0^2 \cdot \mathcal{F}_{00}^2(q^2) + a_0 a_1 \cdot \mathcal{F}_{10}^2(q^2) \right)$$

$$+ a_1^2 \cdot \mathcal{F}_{11}^2(q^2) + c_0^2 \cdot A^2 \mathcal{F}_S^2(q^2)). \quad (24)$$

The last term corresponding to the spin-independent scalar interaction gains coherent enhancement  $A^2$  ( $A$  is the atomic weight of the nucleus in the reaction). The coefficients  $a_{0,1}, c_0$  do not depend on nuclear structure and relate to the parameters  $\mathcal{A}_q, \mathcal{C}_q$  of the effective Lagrangian (21) and to the parameters  $\Delta q, f_s, \hat{f}$  characterizing the nucleon structure. One has the relationships

$$\begin{aligned} a_0 &= (\mathcal{A}_u + \mathcal{A}_d)(\Delta u + \Delta d) + 2\Delta s \mathcal{A}_s, \\ a_1 &= (\mathcal{A}_u - \mathcal{A}_d)(\Delta u - \Delta d), \\ c_0 &= \hat{f} \frac{m_u \mathcal{C}_u + m_d \mathcal{C}_d}{m_u + m_d} + f_s \mathcal{C}_s + \frac{2}{27}(1 - f_s - \hat{f})(\mathcal{C}_c + \mathcal{C}_b + \mathcal{C}_t). \end{aligned} \quad (25)$$

Here  $\Delta q^{p(n)}$  are the fractions of the proton(neutron) spin carried by the quark  $q$ . The standard definition is

$$\langle p(n) | \bar{q} \gamma^\mu \gamma_5 q | p(n) \rangle = 2S_{p(n)}^\mu \Delta q^{p(n)}, \quad (26)$$

where  $S_{p(n)}^\mu = (0, \vec{S}_{p(n)})$  is the 4-spin of the nucleon. The parameters  $\Delta q^{p(n)}$  can be extracted from data on polarized nucleon structure functions [39, 40] and hyperon semileptonic decay data [41].

We use in the analysis  $\Delta q$  values extracted both from the EMC [39] and from SMC [40] data.

The other nuclear structure parameters  $f_s$  and  $\hat{f}$  in formula (25) are defined as follows:

$$\begin{aligned} \langle p(n) | (m_u + m_d)(\bar{u}u + \bar{d}d) | p(n) \rangle &= 2\hat{f} M_{p(n)} \bar{\Psi} \Psi, \\ \langle p(n) | m_s \bar{s}s | p(n) \rangle &= f_s M_{p(n)} \bar{\Psi} \Psi. \end{aligned} \quad (27)$$

The values extracted from the data under certain theoretical assumptions are [43]:

$$\hat{f} = 0.05 \quad \text{and} \quad f_s = 0.14. \quad (28)$$

The strange quark contribution  $f_s$  is known to be uncertain to about a factor of 2. Therefore we take its value in the analysis within the interval  $0.07 < f_s < 0.3$  [42, 43].

The nuclear structure comes into play via the form factors  $\mathcal{F}_{ij}(q^2), \mathcal{F}_S(q^2)$  in Eq. (24). The spin-independent form factor  $\mathcal{F}_S(q^2)$  can be represented as the normalized Fourier transform of a spherical nuclear ground state density

distribution  $\rho(\mathbf{r})$ . In the analysis we use the standard Woods-Saxon inspired distribution [44]. It leads to the form factor

$$\mathcal{F}_S(q^2) = \int d^3\mathbf{r} \rho(\mathbf{r}) e^{i\mathbf{r}\mathbf{q}} = 3 \frac{j_1(qR_0)}{qR_0} e^{-\frac{1}{2}(qs)^2}, \quad (29)$$

where  $R_0 = (R^2 - 5s^2)^{1/2}$  and  $s \approx 1$  fm are the radius and the thickness of a spherical nuclear surface respectively,  $j_1$  is the spherical Bessel function of index 1.

Spin-dependent form factors  $\mathcal{F}_{ij}(q^2)$  are much more nuclear model dependent quantities. The last few years have seen a noticeable progress in detailed nuclear model calculations of these form factors. For many nuclei of interest in DM search they have been calculated within the conventional shell model [45] and within an approach based on the theory of finite Fermi systems [46]. We use the simple parameterization of the  $q^2$  dependence of  $\mathcal{F}_{ij}(q^2)$  in the form of a Gaussian with the *r.m.s.* spin radius of the nucleus calculated in the harmonic well potential [47]. For our purposes this semi-empirical scheme is sufficient.

An experimentally observable quantity is the differential event rate per unit mass of the target material

$$\frac{dR}{dE_r} = \left[ N \frac{\rho_\chi}{m_\chi} \right] \int_{v_{min}}^{v_{max}} dv f(v) v \frac{d\sigma}{dq^2}(v, E_r) \quad (30)$$

Here  $f(v)$  is the velocity distribution of neutralinos in the earth's frame which is usually assumed to be a Maxwellian distribution in the galactic frame.  $v_{max} = v_{esc} \approx 600$  km/s and  $\rho_\chi = 0.3$  GeV·cm<sup>-3</sup> are the escape velocity and the mass density of the relic neutralinos in the solar vicinity;  $v_{min} = (M_A E_r / 2M_{red}^2)^{1/2}$  with  $M_A$  and  $M_{red}$  being the mass of nucleus  $A$  and the reduced mass of the neutralino-nucleus system, respectively. Note that  $q^2 = 2M_A E_r$ .

The differential event rate is the most appropriate quantity for comparing with the observed recoil spectrum and allows one to take properly into account spectral characteristics of a specific detector and to separate the background. However, in many cases the total event rate  $R$  integrated over the whole kinematical domain of the recoil energy is sufficient. It is widely employed in theoretical papers for estimating the prospects for DM detection, ignoring experimental complications which may occur on the way. Notice, that the integrated event rate is less sensitive to details of nuclear structure than the differential one (30). The  $q^2$  shape of the form factors  $\mathcal{F}_{ij}(q^2)$ ,  $\mathcal{F}_S(q^2)$  in Eq. (24) may essentially change from one nuclear model to another. Integration over  $q^2$  as in the case of the total event rate  $R$  reduces this model

dependence. In the present paper we are going to perform a general analysis aimed at searching for domains with extraordinary large values of the event rate  $R$  like those reported in [10]. This is the reason why we use in the analysis the total event rate  $R$ .

## 5 Numerical Analysis

In our numerical analysis we randomly scan the MSSM parameter space within a broad domain

$$1 \text{ GeV} < m_{1/2} < 5 \text{ TeV}, \quad |\mu| < 2 \text{ TeV}, \quad (31)$$

$$1 < \tan\beta < 50, \quad |A_0| < 1 \text{ TeV}, \quad (32)$$

$$0 < m_0 < 5 \text{ TeV}, \quad 50 \text{ GeV} < m_A < 1 \text{ TeV}. \quad (33)$$

In the region where  $\tan\beta \gtrsim 35$  the top Yukawa dominance approximation is not applicable in the RGE. Therefore, we use the procedure developed in [13] which takes into account the bottom and tau Yukawa couplings as well.

The cut-off condition  $R > 0.01$  event/kg/day is implemented in the scanning procedure. It reflects realistic sensitivities of the present and the near-future DM detectors.

Note again, that we use the GUT scenario with the non-universal Higgs mass parameters (see (3)) [11]. The Higgs boson masses are calculated in terms of the CP-odd Higgs boson mass  $m_A$  and other input parameters. If we adopt the ultimate GUT conditions with all scalar mass parameters being equal at the unification scale  $M_X$ , the CP-even Higgs boson mass becomes too big because of strong correlations with the sfermion spectrum. As a result the total event rate  $R$  decreases to small values, typically less than 0.01 event/kg/day.

The main results of our scan are presented in Figs.1-4 in the form of scatter plots. Given in Figs.1-3 are the total event rates  $R$  for  $^{73}\text{Ge}$ ,  $\text{Al}_2\text{O}_3$ , and  $\text{NaI}$  versus neutralino mass  $m_\chi$ , as well as  $R$  versus the ratios  $R_{sd}/R$  of the corresponding spin-dependent ( $R_{sd}$ ) part of  $R$  to  $R$  ( $R = R_{sd} + R_{si}$ ). Fig. 4 presents the neutralino relic density  $\Omega_\chi h^2$  as a function of  $m_\chi$ . All quantities are given with and without the  $b \rightarrow s\gamma$  constraint.

We find that the  $b \rightarrow s\gamma$  constraint strongly reduces the MSSM parameter space. The restriction leaves about 25% of the points of the MSSM parameter space which successfully have passed all other constraints.

This constraint disfavors negative values (in our notation) of the Higgs mixing parameter  $\mu$ . It also shrinks the allowed domain for the parameter  $m_{1/2}$  and consequently reduces the allowed domain for the LSP mass  $m_\chi$ .

As seen from Figs.1-3 the  $b \rightarrow s\gamma$  constraint strongly suppresses those points in the parameter space which correspond to the spin-dominant ( $R_{sd} > R_{si}$ ) event rates in all isotopes analyzed. The observation strengthens the conclusion about dominance of the spin-independent neutralino interaction with nuclei obtained in [11] without  $b \rightarrow s\gamma$  constraint.

Nevertheless it is clear that the large event rates survive the  $b \rightarrow s\gamma$  constraint. In the table 1 we present 5 examples of large event rate points taken from the scatter plots in Fig. 1.

In paper [10] extraordinary large event rates of about 10 events/kg/day for  $^{76}\text{Ge}$  and 100 events/kg/day for NaI were found in a specially arranged scan in the domain  $800 \text{ GeV} < m_\chi < 1200 \text{ GeV}$ ,  $0.01 < Z_g < 0.99$ ,  $0 < m_A < 60 \text{ GeV}$  ( $Z_g = \mathcal{N}_{11}^2 + \mathcal{N}_{12}^2$  see Eq. (12)). We have thoroughly scanned this region to check the cited striking result. We arrived at a negative conclusion. No large event rate domains around  $m_\chi \sim 1 \text{ TeV}$  as quoted in [10] have been found in our scan. Note, since the neutralino is the LSP, these domains correspond to a situation when all SUSY particles are very heavy with masses around 1 TeV or larger. Looking at the formulas (22) and (23) we do not see any natural possibility for  $R$  to approach such large values in this domain. The strong kinematical suppression can only be compensated in the case when  $m_{\tilde{q}} - m_\chi \simeq m_q$ .

## 6 Conclusion

We have systematically studied the allowed MSSM parameter space taking into account various theoretical and experimental constraints. We have found domains with experimentally interesting event rates for the DM neutralino detection ( $R \simeq 10 \text{ events/kg/day}$ ) in the neutralino mass range  $70 \text{ GeV} < m_\chi < 200 \text{ GeV}$ . This would be within the reach of current dark matter experiments. Special attention was paid to the constraint following from the CLEO measurement of  $BR(b \rightarrow s\gamma)$ . We have illustrated that despite the well known fact that this constraint essentially reduces the allowed MSSM parameter space it does not exclude large event rate domains. We have checked the recently reported result [10] on large neutralino detection event rates in the 1 TeV region of the neutralino mass. Our analysis has not reproduced this result.

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### Figure legends

Figure 1. The total event rate  $R$  for  $^{73}\text{Ge}$ , versus mass of neutralino  $m_\chi$  (upper panel) as well as versus the ratio  $R_{sd}/R$  of the corresponding spin-dependent  $R_{sd}$  part of  $R$  to the  $R$  ( $R = R_{sd} + R_{si}$ ). The scatter plots are obtained without (left panel) and with the  $b \rightarrow s\gamma$  constraint (right panel).

Figure 2. The same as in Figure 1, but for sapphire,  $\text{Al}_2\text{O}_3$ .

Figure 3. The same as in Figure 1, but for sodium iodide,  $\text{NaI}$ .

Figure 4. The relic density  $\Omega_\chi h^2$  versus mass of neutralino  $m_\chi$  without (left panel) and with the  $b \rightarrow s\gamma$  constraint (right panel).

**Table**

SUSY points	1	2	3	4	5
$\tan \beta$	20.4	21.2	21.2	12.7	19.5
$m_0$ (GeV)	3654	1421	3055	646	590
$m_{1/2}$ (GeV)	621	229	405	372	320
$A_0$ (GeV)	-2.8	-0.18	0.5	-5.3	-1.4
$m_A$ (GeV)	941	588	673	575	685
$\mu$ (GeV)	176	575	678	606	652
$m_\chi$ (GeV)	70.1	91.3	163	148	129
$\Omega_\chi h_0^2$	0.17	0.074	0.1	0.054	0.16
R(events/kg/day)	7.38	1.08	2.13	2.25	1.73
$R_{sd}/R_{si} \cdot 10^3$	0.5	0.2	0.04	0.07	0.06
Gaugino fraction, $Z_g$	0.04	0.95	0.95	0.93	0.96
$BR(b \rightarrow s\gamma) \cdot 10^3$	0.28	0.22	0.29	0.25	0.19

Table 1. Representative points with large event rate values for germanium,  $^{73}\text{Ge}$ . (The gaugino fraction is defined as  $Z_g = \mathcal{N}_{11}^2 + \mathcal{N}_{12}^2$ )

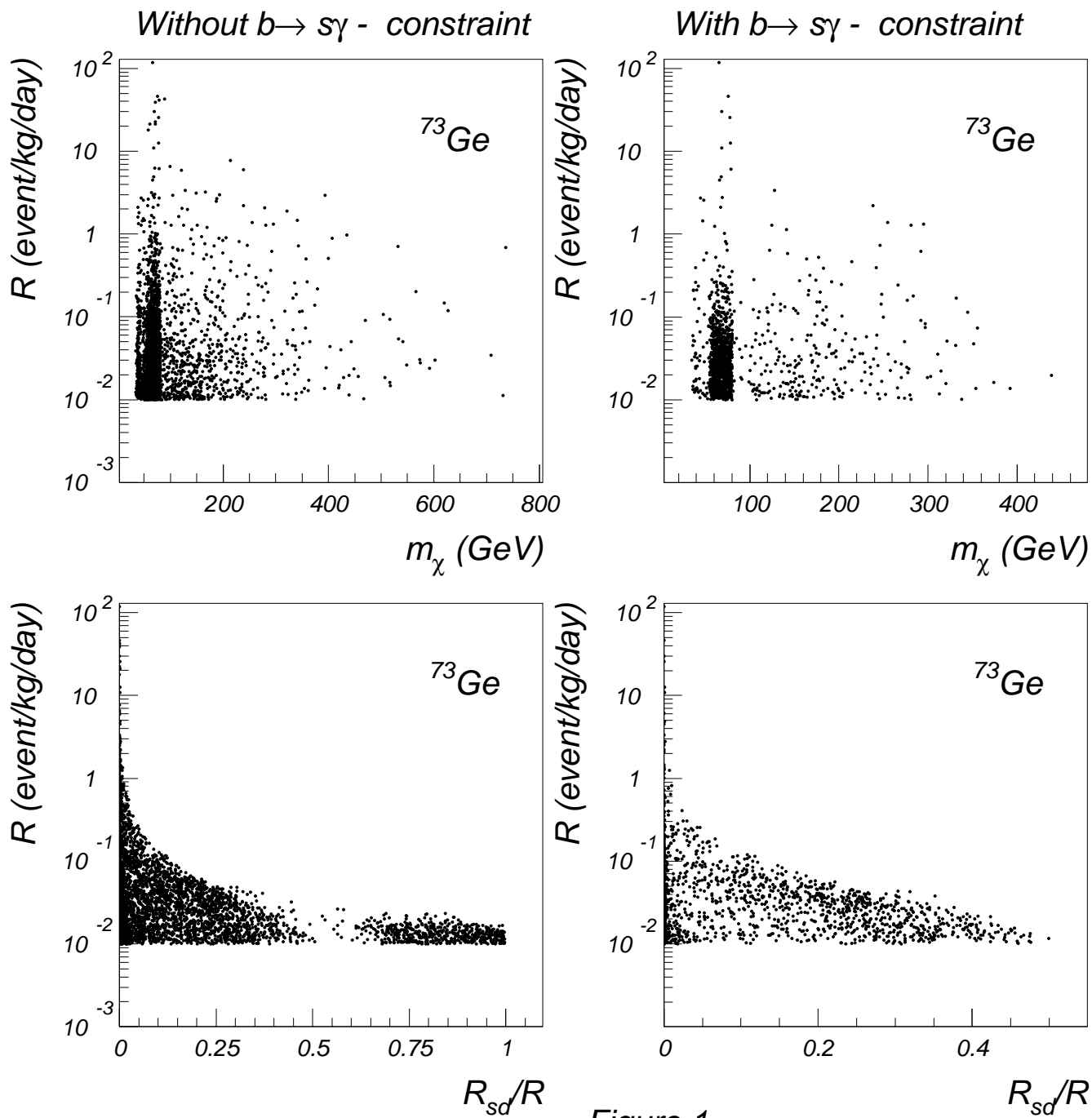


Figure 1

*Without  $b \rightarrow s\gamma$  - constraint*

*With  $b \rightarrow s\gamma$  - constraint*

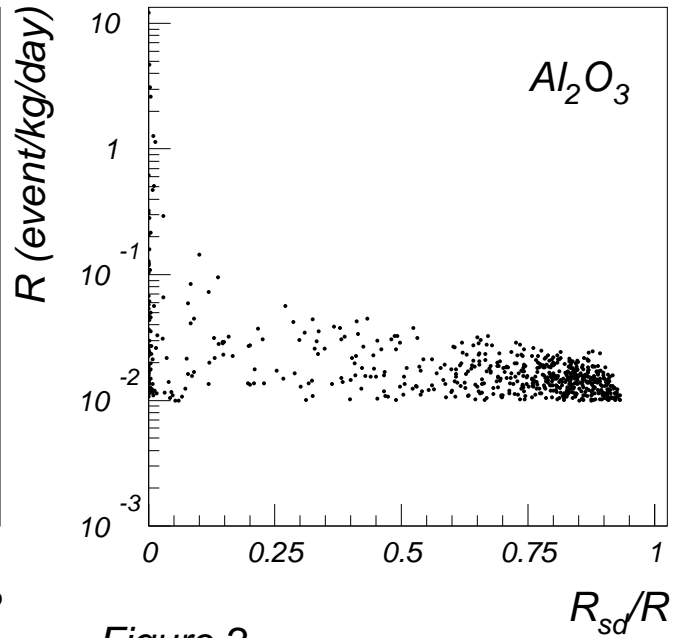
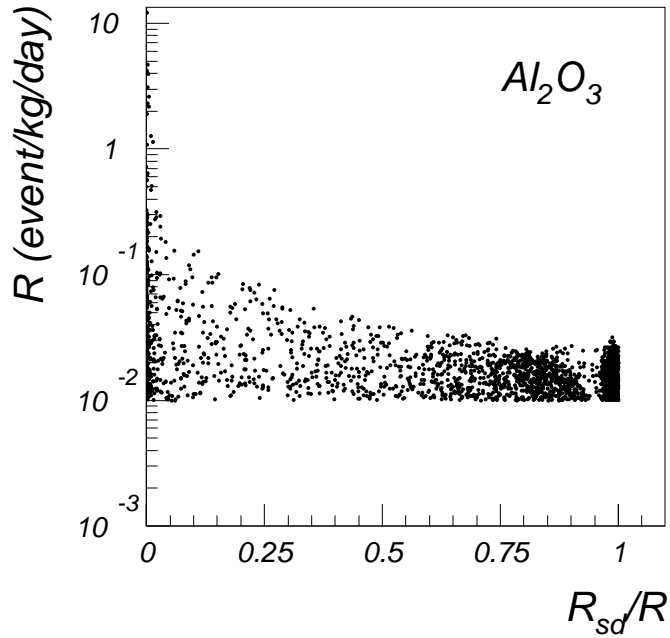
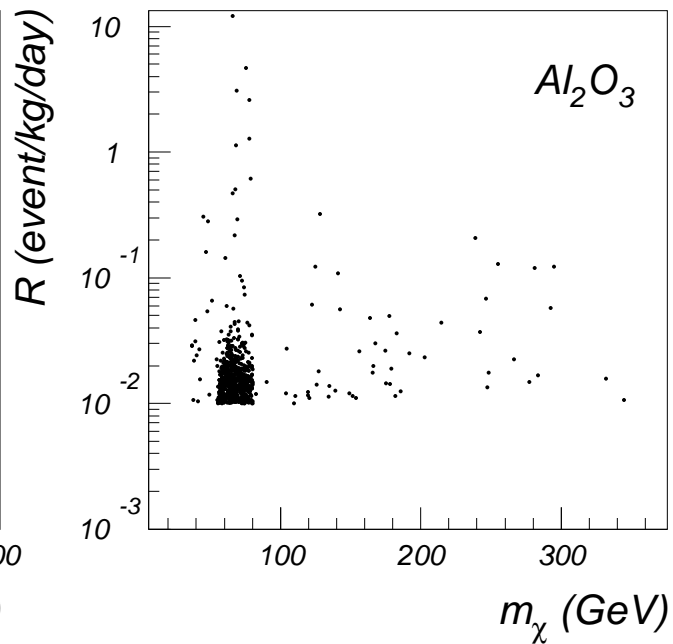
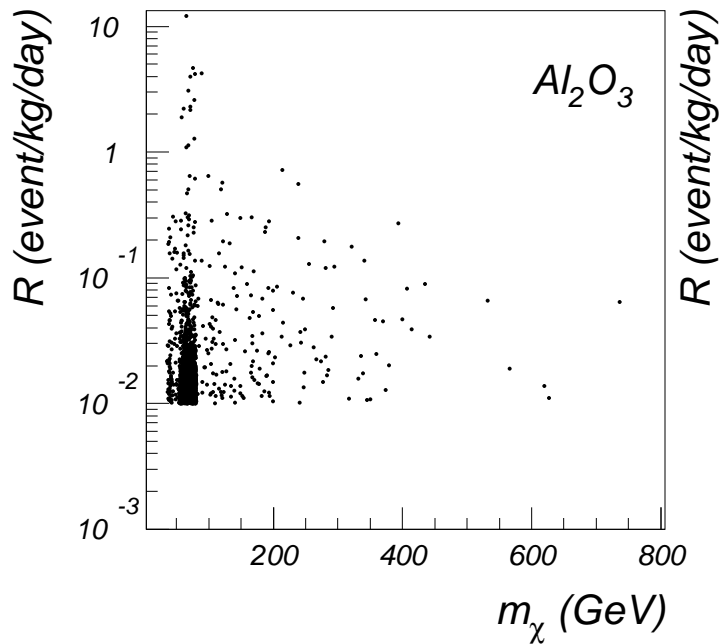


Figure 2

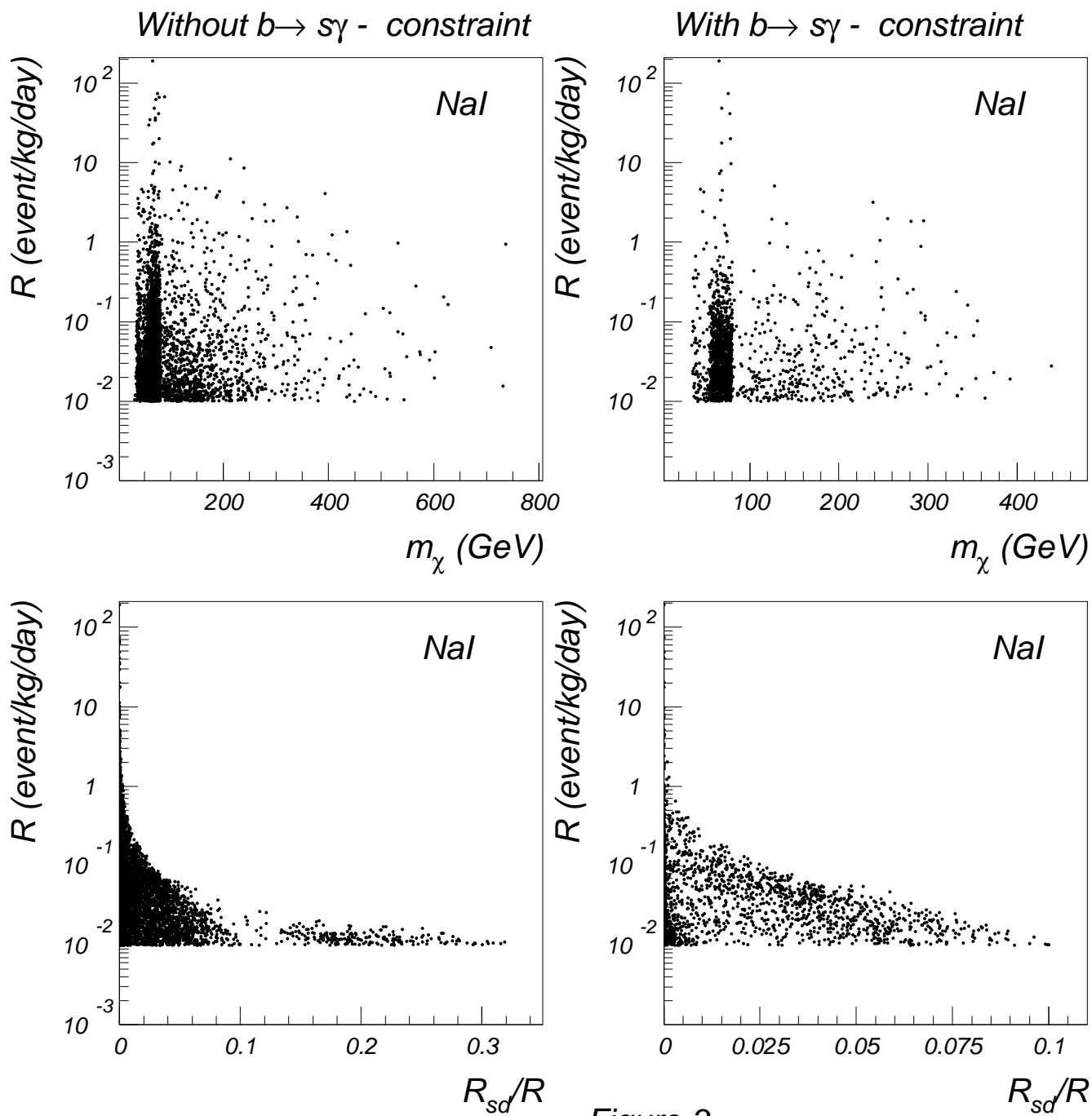


Figure 3



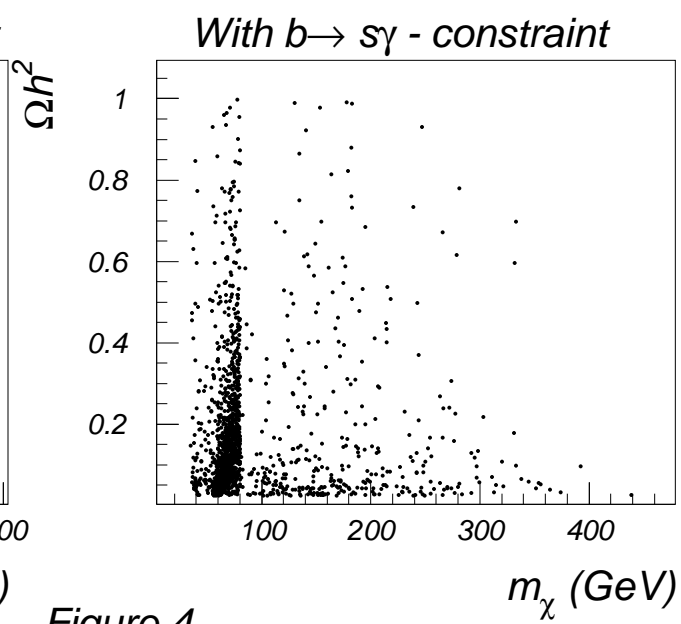
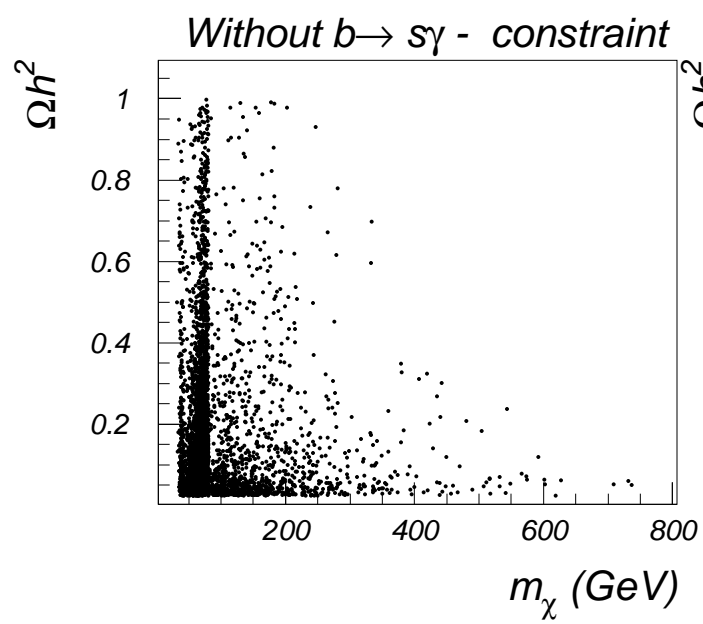


Figure 4